

Is symmetry breaking of $SU(5)$ theory responsible for the diphoton excess?

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We advocate the possibility that the observed diphoton excess at 750 GeV at the LHC can be addressed by the scalar field that is a part of the $SU(5)$ symmetry breaking sector. The field in question is the Standard Model singlet that resides in the adjoint representation that breaks $SU(5)$ down to $SU(3) \times SU(2) \times U(1)$. We also show that the required production and subsequent decay to two photons of this singlet can be induced by individual or combined contribution of two scalar multiplets S_3 and R_2 that transform as $(\mathbf{3}, \mathbf{3}, -1/3)$ and $(\mathbf{3}, \mathbf{2}, 7/6)$ under $SU(3) \times SU(2) \times U(1)$, respectively. The individual dominance of these multiplets is directly related to the issue of the charged fermion mass generation within the $SU(5)$ framework and can be unambiguously tested through the diboson decay signatures of the Standard Model singlet field.

The first results from Run 2 of the LHC experiments have revealed a hint of an unexpected feature in diphoton final state. With integrated luminosity of $\sim 3 \text{ fb}^{-1}$, collected at the center-of-mass energy of 13 TeV, both ATLAS and CMS experiments have reported modest excesses of two-photon events over the Standard Model (SM) background [1, 2]. The global statistical significances are small. However, local significances of the excesses reach 3.9σ and 3.4σ at ATLAS and CMS, respectively. These excesses are furthermore located in the same region of the diphoton invariant mass at $m_{\gamma\gamma} \simeq 750 \text{ GeV}$. The simplest theoretical interpretation of the preliminary diphoton signal is to introduce a scalar particle that is a singlet of the SM and along with it additional fermionic and/or bosonic degrees of freedom that mediate the singlet interaction to pairs of gauge bosons. See Refs. [3–5] for explicit examples.

We advocate the possibility that the SM singlet in question is a part of the $SU(5)$ symmetry breaking sector. Recall, $SU(5)$ is broken down to $SU(3) \times SU(2) \times U(1)$ through a vacuum expectation value (VEV) of the SM singlet field in 24-dimensional scalar representation [6]. The decomposition of the adjoint representation of $SU(5)$ under $SU(3) \times SU(2) \times U(1)$ is $\mathbf{24} \equiv \Sigma = (\mathbf{8}, \mathbf{1}, 0) \oplus (\mathbf{1}, \mathbf{3}, 0) \oplus (\mathbf{3}, \mathbf{2}, -5/6) \oplus (\bar{\mathbf{3}}, \mathbf{2}, 5/6) \oplus (\mathbf{1}, \mathbf{1}, 0) = (\Sigma_8, \Sigma_3, \Sigma_{3,2}, \Sigma_{\bar{3},2}, \Sigma_0)$. Our first goal is to demonstrate that the scalar singlet $\Sigma_0 \equiv (\mathbf{1}, \mathbf{1}, 0)$ can reside at the electroweak scale if needed.

The scalar potential V for Σ is

$$V = -\frac{\mu^2}{2} \Sigma_j^i \Sigma_i^j + \frac{a}{4} (\Sigma_j^i \Sigma_i^j)^2 + \frac{b}{2} \Sigma_j^i \Sigma_k^j \Sigma_l^k \Sigma_l^i + \frac{c}{3} \Sigma_j^i \Sigma_j^j \Sigma_i^i, \quad (1)$$

where μ^2 , a , b , and c represent parameters of the theory. $i, j, k, l, n = 1, \dots, 5$ are the $SU(5)$ indices. We, for definiteness, consider only renormalizable operators. The conditions that the potential V develops a local minimum that breaks the $SU(5)$ down to the SM gauge group are [7]

$$\beta > \begin{cases} \frac{15}{32} \left(\gamma - \frac{4}{15} \right), & \gamma > \frac{2}{15} \\ -\frac{1}{120\gamma}, & \gamma < \frac{2}{15} \end{cases}, \quad (2)$$

where dimensionless variables β and γ are defined as $\beta = (\mu^2 b)/c^2$ and $\gamma = (a/b + 7/15)$, respectively. The symmetry breaking VEV of Σ is $\langle \Sigma \rangle = \lambda/\sqrt{30} \text{ diag}(2, 2, 2, -3, -3)$, where [7]

$$\lambda = \frac{c}{b} \left(\frac{\beta}{\gamma} \right)^{1/2} \left[\left(1 + \frac{1}{120\beta\gamma} \right)^{1/2} + \frac{1}{(120\beta\gamma)^{1/2}} \right] = \frac{c}{b} \left(\frac{\beta}{\gamma} \right)^{1/2} h(\beta\gamma). \quad (3)$$

$\Sigma_{3,2}$ and $\Sigma_{\bar{3},2}$ multiplets are eaten by X and Y gauge bosons of $SU(5)$. These gauge fields mediate proton decay and thus need to be very heavy. Their common mass $m_{(X,Y)}$ is

$$m_{(X,Y)} = \sqrt{\frac{5}{12}} g_{\text{GUT}} \lambda, \quad (4)$$

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where g_{GUT} is the $SU(5)$ gauge coupling at the grand unified theory (GUT) scale m_{GUT} . It is customary to identify $m_{(X,Y)}$ to be the GUT scale, i.e., scale where the SM gauge couplings unify. Potential in Eq. (1) yields the following mass relations

$$\begin{aligned} m_{\Sigma_8}^2 &= \left[\frac{1}{3} + \frac{5}{\sqrt{30}} \left(\frac{\gamma}{\beta} \right)^{1/2} \frac{1}{h(\beta\gamma)} \right] b\lambda^2, \\ m_{\Sigma_3}^2 &= \left[\frac{4}{3} - \frac{5}{\sqrt{30}} \left(\frac{\gamma}{\beta} \right)^{1/2} \frac{1}{h(\beta\gamma)} \right] b\lambda^2, \\ m_{\Sigma_0}^2 &= \left[1 - \frac{1}{1 + (1 + 120\beta\gamma)^{1/2}} \right] 2b\gamma\lambda^2, \end{aligned} \quad (5)$$

where m_{Σ_8} , m_{Σ_3} , and m_{Σ_0} denote masses of Σ_8 , Σ_3 , and Σ_0 , respectively.

We require Σ_0 to be light. We accordingly set $m_{\Sigma_0}^2 = 2b\gamma\lambda^2\epsilon$ to demonstrate viability of this requirement, where $0 < \epsilon \ll 1$. This leads to the following inequality (for $\gamma > 0$)

$$\beta \approx \frac{\epsilon^2 - 1}{120\gamma} > -\frac{1}{120\gamma}. \quad (6)$$

We furthermore obtain $m_{\Sigma_8}^2 = [1/3 + 10\gamma - \mathcal{O}(\epsilon)]b\lambda^2$ and $m_{\Sigma_3}^2 = [4/3 - 10\gamma + \mathcal{O}(\epsilon)]b\lambda^2$. Clearly, the requirement that $m_{\Sigma_3}^2 > 0$ is satisfied for $\gamma < 2/15$. This shows that there exists a part of the parameter space where Σ_0 can reside at the electroweak scale to serve as the candidate behind the diphoton excesses. This possibility is not in collision with the symmetry breaking chain $SU(5) \rightarrow SU(3) \times SU(2) \times U(1)$ at the classical level.

Singlet field Σ_0 of mass $m_{\Sigma_0} \simeq 750 \text{ GeV}$ should couple to vector-like fermions and/or charged scalars in order to be produced at the LHC and to be able to subsequently decay into two photons. Only then will it be able to help explain observed signal excesses [3–5]. Vector-like quarks and leptons are frequently used in GUT model building to address, for example, the issue of the SM fermion masses and mixings. The idea is to mix the SM fermions with one or more of the SM multiplets in these additional $SU(5)$ representations to produce viable masses and mixing parameters [8]. The most commonly used representations to accommodate vector-like states are 5-, 10-, 15-, and 24-dimensional representations. The relevant operators, at the $SU(5)$ level, are straightforward to write down and we omit them in this note. For explicit proposals to couple vector-like representations to an $SU(5)$ singlet to address observed diphoton excess at 750 GeV at the LHC in a non-supersymmetric (supersymmetric) setting see Ref. [9] (Refs. [10, 11]). The use of vector-like multiplets that comprise full $SU(5)$ multiplet(s) of 5- and 10-dimensional nature has also been advocated in Ref. [12].

Scalar states, on the other hand, are necessary since one or more Higgs doublets are needed to generate fermion masses in the first place. Most commonly used scalar representations in $SU(5)$ are accordingly 5- and 45-dimensional ones. The latter representation contains, among other states, two scalar multiplets S_3 and R_2 that transform as $(\mathbf{3}, \mathbf{3}, -1/3)$ and $(\mathbf{3}, \mathbf{2}, 7/6)$ under $SU(3) \times SU(2) \times U(1)$, respectively. (Here we adopt notation of Ref. [13] to denote relevant colored scalar multiplets.) These particular fields can generate required signal strength very efficiently as we show next.

The operators that couple S_3 and R_2 to the singlet Σ_0 , at the $SU(5)$ level, are $m\mathbf{45}_{ij}^{ij}\Sigma_l^k\mathbf{45}_{ij}^{l*}$ and $\sigma\mathbf{45}_{ij}^{ij}\Sigma_l^k\Sigma_n^l\mathbf{45}_{ij}^{n*}$, where m and σ are *a priori* unknown dimensionful and dimensionless coefficients, respectively. We find that the relevant trilinear vertex is $(-m\sqrt{3/10} + 3/5\sigma\lambda)\Sigma_0(S_3^\dagger S_3 + R_2^\dagger R_2)$. The trilinear vertex coefficient is thus the same for both S_3 and R_2 due to the underlying $SU(5)$ symmetry.

From the point of view of an effective theory defined at the 1 TeV scale the most important operator for our phenomenological study is

$$\mathcal{L}^{\text{eff}} \supset x m_{\Sigma_0} \Sigma_0 \left(S_3^\dagger S_3 + R_2^\dagger R_2 \right), \quad (7)$$

where we set $m_{\Sigma_0} = 750 \text{ GeV}$. As shown in the previous paragraph, dimensionless parameter x is directly related to the parameters of the GUT potential. The trilinear vertex of Eq. (7) destabilizes Σ_0 by opening a decay channel to a pair of S_3 's and/or R_2 's, if these are lighter than $m_{\Sigma_0}/2$. We, however, opt to present our analysis in the regime where Σ_0 cannot decay to a pair of on-shell colored scalar states. In such a setting Σ_0 decays predominantly to the pairs of the SM gauge bosons via loops containing electrically charged colored scalars.

We adapt the analogous expressions for the SM Higgs decay widths for $h \rightarrow \gamma\gamma, gg$, in the presence of scalar degrees of freedom, to a particular case of $\Sigma_0 \rightarrow \gamma\gamma, gg$ [14, 15]. (See also Appendix A for more details.) In the case that the R_2 contribution is dominant the width expressions read

$$\begin{aligned} \Gamma(\Sigma_0 \rightarrow \gamma\gamma) &= |x|^2 \frac{\alpha^2 m_{\Sigma_0}^5}{2^{10} \pi^3 m_{R_2}^4} D_{\gamma\gamma}^{R_2} |\mathcal{A}_0(\tau)|^2, \\ \Gamma(\Sigma_0 \rightarrow gg) &= |x|^2 \frac{\alpha_s^2 m_{\Sigma_0}^5}{2^5 \pi^3 m_{R_2}^4} C(R_2)^2 |\mathcal{A}_0(\tau)|^2. \end{aligned} \quad (8)$$

We will also consider regime in which S_3 and R_2 are simultaneously affecting the Σ_0 decays and in order to do that we take into account the decay amplitudes with interference effects included. The charge eigenstates within weak multiplets S_3 and R_2 are assumed to be degenerate with a common masses of m_{S_3} and m_{R_2} , respectively. We denote by α (α_S) the electromagnetic (strong) coupling, the color algebra factor for color triplets is $C(S_3) = 1/2$, whereas $D_{\gamma\gamma}$ represents the boost factor of the diphoton width stemming from the sum over all charge and color eigenstates propagating in the loop. $D_{\gamma\gamma}$ reads

$$D_{\gamma\gamma} = \left\{ d_c(2T+1) \left[Y^2 + \frac{T(T+1)}{3} \right] \right\}^2, \quad (9)$$

where d_c is the dimension of the $SU(3)$ representation of the scalar, Y is its hypercharge and T the weak isospin. Given the strong dependence of $D_{\gamma\gamma}$ on hypercharge and weak isospin it is now evident why we favor at least one of the two scalar triplets with $d_c = 3$, namely S_3 or R_2 , to be light. Their SM quantum numbers — $Y = -1/3, T = 1$ for S_3 and $Y = 7/6, T = 1/2$ for R_2 — yield large diphoton boost factors ($D_{\gamma\gamma}^{S_3} = 49, D_{\gamma\gamma}^{R_2} \approx 93$) compared to majority of other scalars contained within the 45-dimensional representation. For a more vivid comparison, consider colored scalar in the representation $(\mathbf{\bar{3}}, \mathbf{1}, 1/3)$, studied in Ref. [16], or $(\mathbf{3}, \mathbf{2}, 1/6)$ that result in 0.1 and 2.8 for $D_{\gamma\gamma}$, respectively. The loop function of the argument $\tau = m_{\Sigma_0}^2/(4m_{\text{LQ}}^2)$, with $\text{LQ} = S_3, R_2$, reads

$$\mathcal{A}_0(\tau) = \frac{f(\tau) - \tau}{\tau^2}, \quad f(\tau) = \begin{cases} \arcsin^2 \sqrt{\tau} & ; \tau \leq 1, \\ -\frac{1}{4} \left(\log \frac{1+\sqrt{1-1/\tau}}{1-\sqrt{1-1/\tau}} - i\pi \right)^2 & ; \tau > 1 \end{cases}, \quad (10)$$

and is consistent with the decay amplitude expressions we present in Appendix A.

The gluonic decay $\Sigma_0 \rightarrow gg$ dominates the total width Γ_{Σ_0} while the remaining diboson widths are subleading but non-negligible. We accordingly included widths for $\Sigma_0 \rightarrow Z\gamma, ZZ, WW$ processes in the total width Γ_{Σ_0} in our analysis. The ratios of the diboson to diphoton decay widths that we list in Table I, for two cases, where either one or the other of the two colored scalars is dominant, exhibit very little dependence on the colored state mass. Closer inspection of Table I reveals that one could clearly distinguish the two scenarios through the decays of Σ_0 into diboson channels. (Here and in the following we employ for the gauge couplings $\alpha_S(m_{\Sigma_0}/2) = 0.095$ [17] and $\alpha(m_Z) = 0.0078$ [18].)

VV'	$Z\gamma$	ZZ	W^+W^-	gg
$\frac{\Gamma(\Sigma_0 \rightarrow VV')}{\Gamma(\Sigma_0 \rightarrow \gamma\gamma)} \Big _{m_{S_3} \ll m_{R_2}}$	4.3	7.8	26	54
$\frac{\Gamma(\Sigma_0 \rightarrow VV')}{\Gamma(\Sigma_0 \rightarrow \gamma\gamma)} \Big _{m_{R_2} \ll m_{S_3}}$	0.062	0.55	0.85	13
$\frac{\Gamma(\Sigma_0 \rightarrow VV')}{\Gamma(\Sigma_0 \rightarrow \gamma\gamma)} \Big _{m_{R_2} = m_{S_3}}$	0.52	2.6	7.2	27

TABLE I. Ratio of diboson to diphoton decay widths $\Gamma(\Sigma_0 \rightarrow VV')/\Gamma(\Sigma_0 \rightarrow \gamma\gamma)$. The predictions of the S_3 (R_2) dominance case is shown in the first (second) numeric row. Results in the last row are obtained with the assumption of the mass degeneracy for the two colored scalar states.

Several phenomenological analyses revealed the main characterizing feature of the excess observed in $\sigma(pp \rightarrow \gamma\gamma)_{m_{\gamma\gamma} \approx 750 \text{ GeV}}$ at $\sqrt{s} = 13 \text{ TeV}$ (see e.g. [5, 12, 19]). Assuming a narrow scalar diphoton resonance we employ the following value in this work:

$$\sigma(pp \rightarrow \Sigma_0) \text{Br}(\Sigma_0 \rightarrow \gamma\gamma) \approx (3.5 - 7) \text{ fb}. \quad (11)$$

Recorded statistics in ATLAS and CMS datasets are insufficient at the moment to be able to determine the width of the Σ_0 resonance. Good consistency with the dataset is obtained both for large $\Gamma_{\Sigma_0} \lesssim 0.1 m_{\Sigma_0}$, as well as for significantly narrower Γ_{Σ_0} [12, 20, 21]. We can relate the diphoton excess of Eq. (11) with the partial decay widths

$$\sigma(pp \rightarrow \Sigma_0) \text{Br}(\Sigma_0 \rightarrow \gamma\gamma) = \frac{K_{gg} C_{gg}}{m_{\Sigma_0} s} \frac{\Gamma(\Sigma_0 \rightarrow gg) \Gamma(\Sigma_0 \rightarrow \gamma\gamma)}{\Gamma_{\Sigma_0}}, \quad (12)$$

where we employ factor K_{gg} to include higher order QCD corrections, whereas the gluon parton distribution function convolution is embodied in C_{gg} . At center-of-mass energy of the LHC Run 2 of $\sqrt{s} = 13 \text{ TeV}$ we adopt $C_{gg} = 2.1 \times 10^3$ and $K_{gg} \approx 1.5$, where both values are taken from Ref. [5]. We show in the left (right) panel of Fig. 1 the region in $m_{S_3}-x$ ($m_{R_2}-x$) plane that satisfies the constraint given in Eq. (11) assuming that only state S_3 (R_2) is light. The allowed regions presented in Fig. 1 suggest that relatively light colored scalars with coupling x of order 1 can individually accommodate observed signal. From Table I one can observe that other diboson partial widths are enhanced (suppressed) with respect to the diphoton partial width in the S_3 (R_2) dominance scenario. Current searches at the LHC are not yet sensitive to other diboson decays of Σ_0 [5, 22]

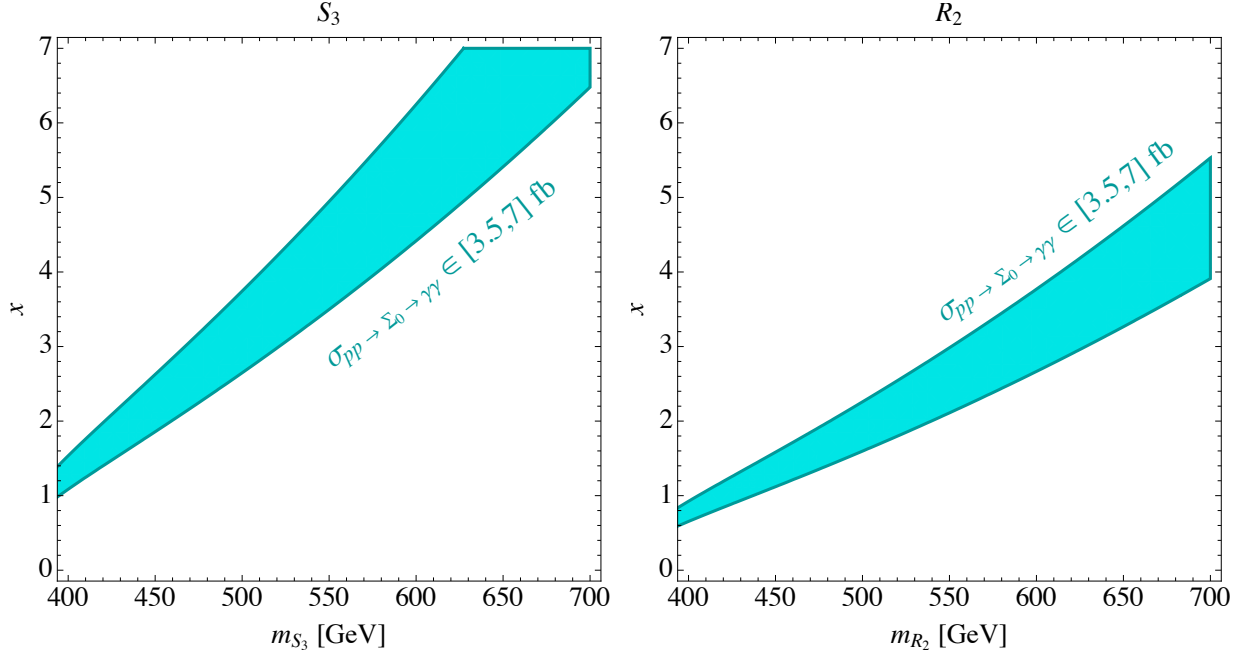


FIG. 1. Left panel: parameter space for the case of S_3 dominance in the m_{S_3} - x plane that satisfies the constraint on $\sigma(pp \rightarrow \Sigma_0 \rightarrow \gamma\gamma) \in [3.5, 7] \text{ fb}$ in cyan. Right panel: parameter space for the case of R_2 dominance.

but this particular feature might be experimentally accessible in near future. This would allow one to probe the nature of the source of the diphoton excess.

It might be the case that both S_3 and R_2 contribute towards the diphoton signal. We accordingly present viable parameter space in the m_{S_3} - m_{R_2} plane for four values of x ($= 0.5, 1, 2, 4$) that yield $\sigma(pp \rightarrow \Sigma_0 \rightarrow \gamma\gamma) \in [3.5, 7] \text{ fb}$ in Fig. 2. We stress again that it is a prediction of $SU(5)$ symmetry that parameter x is the same for both R_2 and S_3 states.

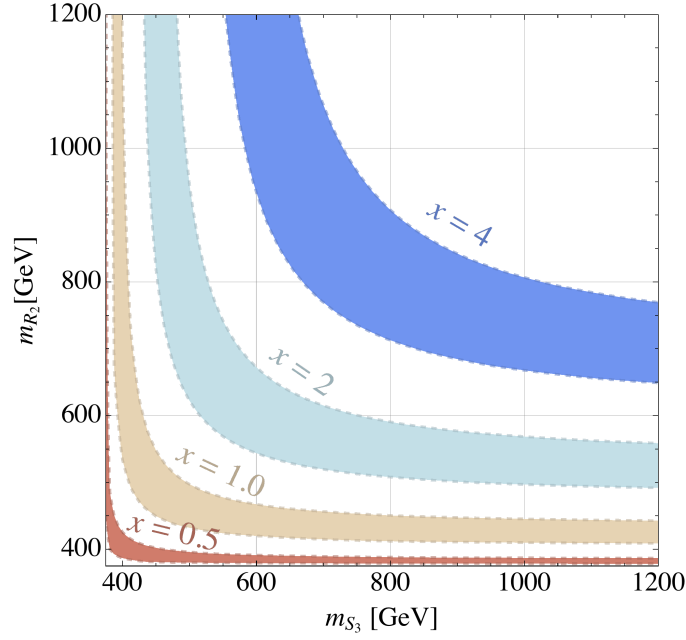


FIG. 2. Parameter space of the scenario that satisfies the constraint on $\sigma(pp \rightarrow \Sigma_0 \rightarrow \gamma\gamma) \in [3.5, 7] \text{ fb}$ in m_{S_3} - m_{R_2} plane for different values of parameter x .

The lightness of either S_3 or R_2 could potentially be in tension with the direct search limits. Both scalars have correct

quantum numbers to be leptoquarks (LQs) [13]. S_3 could, furthermore, mediate proton decay if all possible couplings with the SM fermions are present. We now demonstrate viability of our proposal with a special emphasis on the fermion mass generation within $SU(5)$.

The couplings of S_3 and R_2 in the 45-dimensional representation with the SM fermions that reside in the 10- and 5-dimensional representations of $SU(5)$ originate from two contractions in the $SU(5)$ space. These are $Y_{1\alpha\beta}^{45}\mathbf{10}_\alpha\mathbf{10}_\beta\mathbf{45}$ and $Y_{2\alpha\beta}^{45}\mathbf{10}_\alpha\mathbf{5}_\beta\mathbf{45}^*$, where $\mathbf{10}_\alpha \equiv (\mathbf{1}, \mathbf{1}, 1)_\alpha \oplus (\mathbf{3}, \mathbf{1}, -2/3)_\alpha \oplus (\mathbf{3}, \mathbf{2}, 1/6)_\alpha = (e_\alpha^C, u_\alpha^C, Q_\alpha)$ and $\mathbf{5}_\beta \equiv (\mathbf{1}, \mathbf{2}, -1/2)_\beta \oplus (\mathbf{3}, \mathbf{1}, 1/3)_\beta = (L_\beta, d_\beta^C)$ [6]. The elements of Yukawa coupling matrices are denoted with $Y_{1\alpha\beta}^{45}$ and $Y_{2\alpha\beta}^{45}$, where $\alpha, \beta = 1, 2, 3$ are flavor indices. The $SU(5)$ indices, on the other hand, are suppressed for clarity. One must also introduce one 5-dimensional scalar representation ($\mathbf{5}$) if one wants to generate viable masses of the SM charged fermions through VEVs of electrically neutral Higgs-like fields in $\mathbf{5}$ and $\mathbf{45}$ [23]. The relevant contractions are $Y_{1\alpha\beta}^{45}\mathbf{10}_\alpha\mathbf{10}_\beta\mathbf{5}$ and $Y_{2\alpha\beta}^{45}\mathbf{10}_\alpha\mathbf{5}_\beta\mathbf{5}^*$. We denote VEVs of $\mathbf{5} \equiv \mathbf{5}^i$ and $\mathbf{45} \equiv \mathbf{45}_k^{ij}$ with $\langle \mathbf{5}^5 \rangle = v_5/\sqrt{2}$ and $\langle \mathbf{45}_1^{15} \rangle = \langle \mathbf{45}_2^{25} \rangle = \langle \mathbf{45}_3^{35} \rangle = v_{45}/\sqrt{2}$, where the $SU(5)$ indices are shown for clarity. The mass matrices of the SM charged fermions are

$$m_D = -Y_2^{45}v_{45} - Y_2^5v_5/2, \quad (13)$$

$$m_E = 3Y_2^{45T}v_{45} - Y_2^{5T}v_5/2, \quad (14)$$

$$m_U = 2\sqrt{2}(Y_1^{45} - Y_1^{45T})v_{45} - \sqrt{2}(Y_1^5 + Y_1^{5T})v_5, \quad (15)$$

where the VEVs are taken to be real. The VEV normalization yields $v_5^2/2 + 12v_{45}^2 = v^2$, where $v (= 246 \text{ GeV})$ is the electroweak VEV [24]. m_D , m_E , and m_U are 3×3 mass matrices for down-type quarks, up-type quarks, and charged leptons in flavor basis, respectively.

Let us first assume that the only operators present are the ones that are needed to generate viable masses of the SM charged fermions [23]. These operators are proportional to Y_1^5 , Y_2^5 , and Y_2^{45} . The couplings of both R_2 and S_3 that originate from contraction $Y_{2\alpha\beta}^{45}\mathbf{10}_\alpha\mathbf{5}_\beta\mathbf{45}^*$ are of the leptoquark nature. (See Table II of Ref. [25] for explicit evaluation of the aforementioned contraction with regard to the S_3 couplings.) Note that S_3 can only couple with the quark doublet and the leptonic doublet in this particular instance in the gauge invariant way. Moreover, the requirement to have experimentally viable masses for the SM fermions predicts prompt decays of S_3 and R_2 if and when these are produced at the LHC. To demonstrate that prediction it is sufficient to eliminate Y_2^5 from Eqs. (13) and (14). The GUT scale relation $m_E^T - m_D = 4Y_2^{45}v_{45}$ and the fact that $2m_b(m_{\text{GUT}}) \approx m_\tau(m_{\text{GUT}}) = 1.56 \text{ GeV}$ [24] enable one to establish a lower bound on the largest matrix element in $|Y_2^{45}|$. If we take the limit $v_5 \rightarrow 0$ we find that the largest entry of $|Y_2^{45}|$ exceeds 2×10^{-4} . ($m_b(m_{\text{GUT}})$ and $m_\tau(m_{\text{GUT}})$ are masses of bottom quark and τ lepton at the GUT scale, respectively. In non-supersymmetric setting the running of these masses does not depend strongly on the threshold corrections.) Note that this is conservative bound since v_5 cannot be too small in order to produce top quark mass through Eq. (15). This result implies that the direct searches for the scalar LQ states at the LHC are applicable to this particular scenario.

The lower mass limits on R_2 and S_3 within this particular *ansatz* thus originate from three complementary types of experimental searches for leptoquarks at the LHC. The most stringent limits originate from (i) a search for pair production of first generation LQs [26, 27], (ii) a search for pair production of second generation LQs [26, 27], and (iii) a search for pair production of third generation scalar LQs [28, 29]. The most relevant bounds from these searches are all based on the data sets collected at the LHC in proton-proton collisions at the center-of-mass energy of $\sqrt{s} = 8 \text{ TeV}$. The most constraining of these lower bounds is the one on the mass of second generation LQs that is at 1080 GeV, where the branching fraction of LQ to decay into a charged lepton–quark pair is taken to be equal to one. We will concentrate on R_2 and argue that m_{R_2} can actually be as low as 400 GeV and still be experimentally allowed. For the other leptoquark we will take that the most conservative experimental bound applies to simplify discussion.

The operator $Y_{2\alpha\beta}^{45}\mathbf{10}_\alpha\mathbf{5}_\beta\mathbf{45}^*$ implies that R_2 couples to the right-handed up-type quarks and the leptonic doublets. More specifically, the R_2 component with the electric charge of 5/3 (2/3) couples to the right-handed up-type quarks and charged (neutral) leptons. Let us explicitly assume that the R_2 component with the 5/3 charge decays 50% of the time into the top– τ lepton pairs and 50% of the time into the top– μ pairs. The R_2 component with 2/3 charge then decays 100% into a top– ν final state. (One needs to perform summation over ν 's in the final state.) There are no stringent constraints from the LHC on the top– ν decays for the pair production of leptoquarks. And, the fact that branching fraction is 0.5 for going into the top– τ lepton pairs tells us that the LHC bound on the R_2 component with the 5/3 charge is roughly 500 GeV. See right panel of Fig. 3 in Ref. [30] for the implications of the recast of the search for LQs that decay into the top– τ lepton pairs that was performed by the CMS Collaboration [29]. We can make the 5/3 component of R_2 as light as 400 GeV if needed by allowing additional decays into the top– e pairs. Note that it is reasonable to assume that the entries of Y_2^{45} that are associated with the third generation dominate. In our case we assume that the third row of Y_2^{45} is the dominant one to demonstrate that R_2 can be very light. This *ansatz* is also stable under the renormalization group equation (RGE) running from the GUT scale down to the electroweak scale. The stability under the RGE running when one column (or row) in the Yukawa matrix that determines the LQ couplings to the SM fermions dominates has been demonstrated in Ref. [31].

In view of the preceding discussion we choose to vary m_{R_2} from 400 to 700 GeV in the right panel of Fig. 1. (Scalar LQ

multiplets that transform differently with regard to the SM gauge group have also been proposed to help accommodate diphoton excess in Refs. [16, 32]. Viable scenarios with vector LQs have been presented in Refs. [33, 34].)

The second case we discuss is that the only operators present are the ones proportional to Y_1^5 , Y_2^5 , and Y_1^{45} . The most important feature of the contraction $Y_{1\alpha\beta}^{45} \mathbf{10}_\alpha \mathbf{10}_\beta \mathbf{45}$ is that S_3 has only the “diquark” couplings with the SM fermions in the 10-dimensional representation. This is easy to understand since an $SU(2)$ triplet like S_3 should couple to a pair of the $SU(2)$ quark doublets in order to create an invariant operator under $SU(3) \times SU(2) \times U(1)$ since the leptonic doublet is not at the disposal. (See Table II of Ref. [25] for explicit evaluation of the aforementioned contraction with regard to the S_3 couplings.) $SU(5)$ gauge group also dictates that the S_3 couplings to quarks are antisymmetric in flavor space. We accordingly require that Y_1^{45} is not a symmetric matrix in order to insure that S_3 is coupled to quarks. Note that the issue of the mismatch between the down-type quarks and the charged leptons is not addressed in this particular instance. To do that one would need to introduce, for example, additional vector-like representations.

The most current constraints that are relevant for the allowed mass of the S_3 multiplet components, if S_3 is of “diquark” nature, originate from a search for pair-produced resonances decaying to jet pairs in pp collisions at the LHC [35]. We conservatively interpret these measurements to imply lower limit on the mass of the S_3 “diquark” to be at 390 GeV. This is thus adopted as the lowest value of parameter m_{S_3} that we use to present our results in the left panel of Fig. 1. This time around R_2 couples to the right-handed charged leptons and the quark doublets. More specifically, the R_2 component with the 5/3 (2/3) charge couples to the right-handed charged leptons and up-type (down-type) quarks. In this instance the most conservative experimental limit for the mass of R_2 leptoquark holds true.

The last scenario, and the most general one, that we want to address is when all four operators that contribute towards charged fermion masses are present. In this case S_3 has both “diquark” and leptoquark couplings [25]. This simply means that the proton decay constraints stipulate that S_3 cannot contribute towards the diphoton signal. This particular scenario corresponds to the R_2 dominance that is shown in the right panel of Fig. 1 and the predictions in the third row of Table I. Note that it is sufficient that the entries of Y_2^{45} dominate over entries in Y_1^{45} , where Y_2^{45} has a form that predominantly couples the R_2 component with the 5/3 charge to the right-handed top quark and charged leptons.

Let us finally address the issue of unification of gauge couplings within the non-supersymmetric $SU(5)$ framework with 5- and 45-dimensional scalar representations. To do that we first define quantities $b_{ij}^J = (b_i^J - b_j^J)$, $i, j = 1, 2, 3$, where b_i^J are the β -function coefficients of particle J with mass m_J . b_1^J , b_2^J , and b_3^J are associated with $U(1)$, $SU(2)$, and $SU(3)$ of the SM, respectively. We furthermore introduce coefficients $B_{ij} = \sum_J b_{ij}^J r_J$, where the sum goes through all particles that reside below the GUT scale and parameter $r_J = (\ln m_{\text{GUT}}/m_J)/(\ln m_{\text{GUT}}/m_Z)$ describes where between Z boson mass and the GUT scale particle J is.

The gauge coupling at the GUT scale α_{GUT} is well-behaved in non-supersymmetric $SU(5)$ framework and it can, accordingly, be eliminated using three equations that describe running of individual gauge couplings below the GUT scale. This leaves two relevant equations that read [36]

$$\frac{B_{23}}{B_{12}} = \frac{5 \sin^2 \theta_W - \alpha/\alpha_S}{8 \cdot 3/8 - \sin^2 \theta_W} = 0.721 \pm 0.004, \quad (16)$$

$$\ln \frac{m_{\text{GUT}}}{m_Z} = \frac{16\pi}{5\alpha} \frac{3/8 - \sin^2 \theta_W}{B_{12}} = \frac{184.8 \pm 0.1}{B_{12}}, \quad (17)$$

where we use $\alpha_S(m_Z) = 0.1193 \pm 0.0016$, $\alpha^{-1}(m_Z) = 127.906 \pm 0.019$, and $\sin^2 \theta_W = 0.23126 \pm 0.00005$ [18] to produce numerical values in the right-hand sides of Eqs. (16) and (17). Eq. (16), if satisfied, insures that the gauge couplings meet whereas Eq. (17) provides the corresponding value of the GUT scale.

The SM content yields $B_{23}^{\text{SM}}/B_{12}^{\text{SM}} = 0.53$ instead of the experimentally required value given in Eq. (16). Ideally, one would like to have a light field J with positive b_{23}^J and negative b_{12}^J . This would not only help in bringing the left-hand side of Eq. (16) in agreement with the required experimental value but would also raise the GUT scale m_{GUT} through Eq. (17). As it turns out, S_3 is an ideal candidate with $b_{23}^{S_3} = 9/6$ and $b_{12}^{S_3} = -27/15$. The corresponding coefficients of leptoquark R_2 are $b_{23}^{R_2} = 1/6$ and $b_{12}^{R_2} = 17/15$. We find that unification is possible for light S_3 and R_2 . For example, if we set $m_{S_3} = 400$ GeV and $m_{R_2} = 2$ TeV we obtain exact unification for central values of input parameters with an upper bound on the GUT scale that is $m_{\text{GUT}} \leq 6 \times 10^{15}$ GeV. The particle content comprises three scalar representations, i.e., 5-, 24-, and 45-dimensional representations of $SU(5)$, one adjoint representation with the gauge fields, and the SM fermions. Unification is obtained assuming that all proton decay mediating scalars are at or above 10^{12} GeV. (The maximal value of the GUT scale grows with the mass of R_2 leptoquark.) This demonstrates that light S_3 and/or R_2 represent viable options within non-supersymmetric $SU(5)$ framework with 5- and 45-dimensional scalar representations.

Our proposal opens up a possibility to have one light SM singlet at the electroweak scale in practically any $SU(5)$ setting without the need to introduce *ad hoc* scalars. We furthermore demonstrate viability of our proposal using individual or combined contributions towards diphoton signal of scalar multiplets that transform as $(\mathbf{3}, \mathbf{3}, -1/3)$ and $(\mathbf{3}, \mathbf{2}, 7/6)$ under the SM gauge group. We relate the existence of these colored scalars to the issue of fermion mass generation in $SU(5)$ and provide predictions for the diboson decays of the scalar singlet state at the LHC.

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Appendix A: Diboson decay amplitudes

The decay amplitude of a scalar resonance to diboson final states, $\Sigma_0(q) \rightarrow V(p, \epsilon)V'(p', \epsilon')$, can be expressed in terms of two form factors

$$\mathcal{A}_{\Sigma_0 \rightarrow VV'} = \frac{-im_{\Sigma_0}}{2\pi} \left[A_{VV'} g^{\mu\nu} - 2B_{VV'} \frac{p'^\mu p^\nu}{m_{\Sigma_0}^2} \right] \epsilon_\mu^* \epsilon_\nu'^* . \quad (\text{A1})$$

Ward identity states that the amplitude (A1) vanishes whenever we replace external polarization of a photon or a gluon with its momentum, and this requires that form factors $A_{VV'}$ and $B_{VV'}$ are not independent. Notice that transversality conditions, $\epsilon \cdot p = \epsilon' \cdot p' = 0$, allow replacing $p'^\mu p^\nu$ by $q^\mu q^\nu$ in Eq. (A1), however, in this case one has to enforce transversality also in the polarization sum prescription: $\sum_\lambda \epsilon_\mu(p, \lambda) \epsilon_\nu^*(p, \lambda) \rightarrow -g_{\mu\nu} + \frac{p_\mu p_\nu}{p^2}$, regardless of whether vector $\epsilon(p, \lambda)$ is massless or not. For each diboson decay amplitude mediated by the S_3 state the form factors $A_{VV'}$, $B_{VV'}$, that we present below, have been reduced to the Passarino-Veltman functions with the help of FeynCalc [37, 38] and numerically evaluated using the LoopTools package [39]. In the following expressions, gluon indices are denoted by $A, B (= 1, \dots, 8)$, and one has to insert the quantum numbers of S_3 , i.e., $T = 1, Y = -1/3$. Weak mixing factors $\tan \theta_W$ and $\sin \theta_W$ ($\sin^2 \theta_W = 0.231$) are abbreviated as t_θ and s_θ , respectively. The amplitudes for the cases that involve a massless boson in the final state read:

$$A_{g^A g^B} = B_{g^A g^B} = x\alpha_S \frac{\delta^{AB}}{2} (2T+1) \left[1 + 2m_{\Sigma_0}^2 C_0(0, m_{\Sigma_0}^2, 0, m_{S_3}^2, m_{S_3}^2, m_{S_3}^2) \right], \quad (\text{A2})$$

$$A_{\gamma\gamma} = B_{\gamma\gamma} = x\alpha N_c (2T+1) \left[Y^2 + \frac{T(T+1)}{3} \right] \left[1 + 2m_{\Sigma_0}^2 C_0(0, m_{\Sigma_0}^2, 0, m_{S_3}^2, m_{S_3}^2, m_{S_3}^2) \right], \quad (\text{A3})$$

$$\begin{aligned} A_{Z\gamma} &= \frac{B_{Z\gamma}}{1 - \frac{m_Z^2}{m_{\Sigma_0}^2}} \\ &= x \frac{\alpha N_c}{t_\theta} (2T+1) \left[-Y^2 t_\theta^2 + \frac{T(T+1)}{3} \right] \left\{ 1 + \frac{m_Z^2 [B_0(m_Z^2, m_{S_3}^2, m_{S_3}^2) - B_0(m_{\Sigma_0}^2, m_{S_3}^2, m_{S_3}^2)]}{m_Z^2 - m_{\Sigma_0}^2} \right. \\ &\quad \left. + m_{S_3}^2 [C_0(0, m_{\Sigma_0}^2, m_Z^2, m_{S_3}^2, m_{S_3}^2, m_{S_3}^2) + C_0(m_Z^2, m_{\Sigma_0}^2, 0, m_{S_3}^2, m_{S_3}^2, m_{S_3}^2)] \right\}, \end{aligned} \quad (\text{A4})$$

whereas for the massive final states one finds:

$$A_{ZZ} = x \frac{\alpha N_c}{t_\theta^2} (2T+1) \left[Y^2 t_\theta^4 + \frac{T(T+1)}{3} \right] \quad (\text{A5})$$

$$\times \left\{ 1 + \frac{2m_Z^2 [B_0(m_{\Sigma_0}^2, m_{S_3}^2, m_{S_3}^2) - B_0(m_Z^2, m_{S_3}^2, m_{S_3}^2)]}{m_{\Sigma_0}^2 - 4m_Z^2} \right. \\ \left. + 2 \left[m_{S_3}^2 + \frac{m_Z^4}{m_{\Sigma_0}^2 - 4m_Z^2} \right] C_0(m_Z^2, m_{\Sigma_0}^2, m_Z^2, m_{S_3}^2, m_{S_3}^2, m_{S_3}^2) \right\},$$

$$B_{ZZ} = x \frac{\alpha N_c}{t_\theta^2} (2T+1) \left[Y^2 t_\theta^4 + \frac{T(T+1)}{3} \right] \frac{1}{m_{\Sigma_0}^2 - 4m_Z^2} \quad (\text{A6})$$

$$\times \left\{ m_{\Sigma_0}^2 - 2m_Z^2 - \frac{2m_Z^2 (m_{\Sigma_0}^2 + 2m_Z^2) [B_0(m_Z^2, m_{S_3}^2, m_{S_3}^2) - B_0(m_{\Sigma_0}^2, m_{S_3}^2, m_{S_3}^2)]}{m_{\Sigma_0}^2 - 4m_Z^2} \right. \\ \left. + 2 \left[m_{S_3}^2 (m_{\Sigma_0}^2 - 2m_Z^2) + 2m_Z^4 \left(1 + \frac{3m_Z^2}{m_{\Sigma_0}^2 - 4m_Z^2} \right) \right] C_0(m_Z^2, m_{\Sigma_0}^2, m_Z^2, m_{S_3}^2, m_{S_3}^2, m_{S_3}^2) \right\},$$

$$A_{WW} = x \frac{2\alpha N_c}{s_\theta^2} \left\{ 1 + \frac{2m_W^2 [B_0(m_{\Sigma_0}^2, m_{S_3}^2, m_{S_3}^2) - B_0(m_W^2, m_{S_3}^2, m_{S_3}^2)]}{m_{\Sigma_0}^2 - 4m_W^2} \right. \quad (\text{A7})$$

$$\left. + 2 \left(m_{S_3}^2 + \frac{m_W^4}{m_{\Sigma_0}^2 - 4m_W^2} \right) C_0(m_W^2, m_{\Sigma_0}^2, m_W^2, m_{S_3}^2, m_{S_3}^2, m_{S_3}^2) \right\},$$

$$B_{WW} = x \frac{2\alpha N_c}{s_\theta^2} \frac{1}{(m_{\Sigma_0}^2 - 4m_W^2)} \quad (\text{A8})$$

$$\times \left\{ m_{\Sigma_0}^2 - 2m_W^2 - \frac{2m_W^2 (m_{\Sigma_0}^2 + 2m_W^2) [B_0(m_W^2, m_{S_3}^2, m_{S_3}^2) - B_0(m_{\Sigma_0}^2, m_{S_3}^2, m_{S_3}^2)]}{m_{\Sigma_0}^2 - 4m_W^2} \right. \\ \left. + 2 \left[m_{S_3}^2 (m_{\Sigma_0}^2 - 2m_W^2) + \frac{2m_W^4 (m_{\Sigma_0}^2 - m_W^2)}{m_{\Sigma_0}^2 - 4m_W^2} \right] C_0(m_W^2, m_{\Sigma_0}^2, m_W^2, m_{S_3}^2, m_{S_3}^2, m_{S_3}^2) \right\}.$$

The amplitudes due to virtual R_2 contributions are obtained by adjusting the mass $m_{S_3} \rightarrow m_{R_2}$, inserting appropriate values for T and Y for the electrically neutral final state amplitudes, and adjusting the WW amplitudes by a factor of $1/4$.

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